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GLUINO DECAYS AND EXPERIMENTAL SIGNATURES

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Supersymmetry could be revealed by the discovery of gluinos, the supersymmetric partners of the gluons. Previously, it has been assumed that gluinos decay entirely into $q\bar{q}\tilde{\gamma}$ (or $g\tilde{G}$ if a goldstino $\tilde{G}$ exists) where the photino $\tilde{\gamma}$ is assumed to be the lightest supersymmetric particle. We point out that the C-violating decay $\tilde{g} \to g\tilde{\gamma}$ could be competitive if there is sufficient mass splitting between the scalar partners of the left- and right-handed quarks. We also consider four-body gluino decay modes which under certain conditions could dominate, and we relax the usual assumption of massless photinos. We discuss the phenomenological implications of these new modes and indicate the effects on possible experimental signatures of supersymmetry; when the $g\tilde{\gamma}$ mode is large it greatly improves the possibilities for detecting gluinos in both collider and beam dump experiments, perhaps even allowing detection of individual events at colliders.

1. Introduction

Currently, there is no experimental evidence either positive or negative for the existence of supersymmetry. It is thus extremely important to carefully study (in as model independent a way as possible) supersymmetric theories to uncover their most likely phenomenological signatures [1]. This will allow experimentalists to obtain tighter bounds on the masses of the various supersymmetric particles and perhaps help to eventually uncover their existence.

In this paper, we concentrate on the gluinos ($\tilde{g}$), the supersymmetric partner of the gluon (g). Although the mass of the gluino is a priori an arbitrary parameter, the coupling of gluinos to matter is precisely determined (though unknown scalar quark

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masses affect rate estimates). In particular, the $\tilde{g}\tilde{g}g$ vertex occurs with strength g_s (the strong QCD couplings constant), which suggests that gluinos will be copiously produced at hadronic colliders [2,3] if they exist. Presumably, gluinos would be constituents of new supersymmetric hadrons which would then decay into ordinary matter and a lighter supersymmetric particle [2–6]. Thus detection of gluinos depends on understanding their possible decay modes. Careful study is especially important for gluinos as their main signature at colliders is missing p_T without hard charged leptons, so large numbers of events could have occurred in existing data without being noticed.

In sect. 2, we review the "standard" gluino decay modes previously discussed in the literature, namely $\tilde{g} \to q\bar{q}\tilde{\gamma}$ (and $\tilde{g} \to g\tilde{G}$ if the goldstino $\tilde{G}$ exists) [2,3,5] where the photino $\tilde{\gamma}$ is assumed to be the lightest supersymmetric particle. These decays imply certain experimental signatures of gluinos which we summarize. We then compare these "standard" decay modes with two modes which have not previously been considered. First is the charge conjugation (C) violating decay $\tilde{g} \to g\tilde{\gamma}$. We show that this mode could be competitive with $q\bar{q}\tilde{\gamma}$ if there is sufficient mass splitting between the scalar partners ($\tilde{q}_L, \tilde{q}_R$) of the left- and right-handed quarks (for any flavor). Since the two-body mode will have a much cleaner signature than decays with three or more particles, the existence of the $g\tilde{\gamma}$ mode could greatly improve detection possibilities. If individual gluino production events had a clear signature, rather than requiring an effect in distributions as is the situation when $\tilde{g} \to q\bar{q}\tilde{\gamma}$, the gluino mass to which collider experiments are sensitive would be greatly increased, perhaps doubled. Second is a class of decay modes which could be relevant if the scalar neutrino ($\tilde{\nu}$) were the lightest supersymmetric particle [7,8]. In this case, four-body decays such as $\tilde{g} \to q\bar{q}\nu\tilde{\nu}$ or $\tilde{g} \to q\bar{q}\ell\tilde{\nu}$ are possible and may dominate.

In sect. 3, we provide some of the details for the computation of the $\tilde{g} \to g\tilde{\gamma}$ decay rate. In sect. 4, we examine the experimental signatures if the $\tilde{g} \to g\tilde{\gamma}$ mode is significant. In sect. 5, we compute the four-body decay rates and discuss the corresponding experimental signals for such modes. Our summary is contained in sect. 6, and some relevant mathematical information is relegated to an appendix.

2. Decay modes of the gluino

First we summarize the decay modes of the gluino previously discussed in refs. [2,3,5]. It is necessary to distinguish between two classes of supersymmetric models. Consider first models of spontaneously broken global supersymmetry. Such models contain one exactly massless fermion $\tilde{G}$ called the goldstino [9]. Current algebra arguments [10] allow one to compute the couplings of $\tilde{G}$; the strength of the coupling varies inversely with the scale M_s at which supersymmetry breaks. The

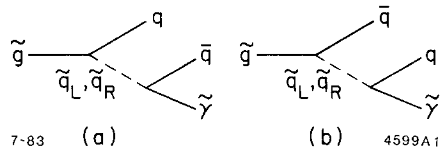

Fig. 1. Three-body decays of the gluino. Because the gluino is a Majorana fermion, we must include both the final states (a) and their charged conjugates, (b). Solid lines are fermions and dashed lines are scalar quarks. The quarks q may run over all possible flavors such that the decay is kinematically allowed.

result for the lifetime $\tilde{g} \to g\tilde{G}$ is then:

$$\tau = \frac{8\pi M_s^4}{\tilde{M}_g^5} = 1.65 \times 10^{-23} \times M_s^4 \tilde{M}_g^{-5} \text{ seconds}, \tag{1}$$

where all masses are given in GeV ($\tilde{M}_g$ = gluino mass). If $M_s \lesssim 1$ TeV the $g\tilde{G}$ mode would dominate if the scalar quarks were heavier than about 100 GeV [3]. However, the current prejudice is to take M_s much larger: $M_s \sim (M_P m_W)^{1/2} \sim 10^{10}$ GeV (where M_P is either the grand unification scale [11] or the Planck mass [12]). For such a large value of M_s, the decay rate, eq. (1), becomes negligible. A second class of models has either softly broken global supersymmetry [13] or low-energy broken supergravity [12, 14, 15]. In the former case, there is no goldstino and in the latter case, the goldstino becomes the helicity $\pm\frac{1}{2}$ component of a massive gravitino $g_{3/2}$. Typically, models of low-energy broken supergravity have $m^2_{g_{3/2}} \approx (0.1\text{–}10) m_W^2$ in which case the decay $\tilde{g} \to g g_{3/2}$ is either kinematically forbidden or irrelevant (since apart from the phase space factor, it is still appropriate to use eq. (1) with $M_s = 10^{10}$ GeV). Thus all currently fashionable models do not have a significant two-body decay $\tilde{g} \to g\tilde{G}$.

We therefore turn to the most probable dominant decay mode $\tilde{g} \to \tilde{\gamma} + q\bar{q}$*. The diagrams which contribute are given in fig. 1. Since $\tilde{g}$ is a (self-conjugate) Majorana fermion, we must sum over all possible final states (fig. 1a) and their charge conjugates (fig. 1b). For simplicity, we will assume that the scalar quarks $\tilde{q}_L$ and $\tilde{q}_R$ (with masses $\tilde{m}_L$ and $\tilde{m}_R$ respectively) are the appropriate scalar mass eigenstates. Then, the appropriate supersymmetric interaction terms are:

$$e_q e\sqrt{2}\left[\bar{\tilde{\gamma}}(P_L \tilde{q}_L^* - P_R \tilde{q}_R^*) q + \bar{q}(P_R \tilde{q}_L - P_L \tilde{q}_R)\tilde{\gamma}\right], \tag{2}$$

where P_R (P_L) is the right- (left-) handed projection operator $\frac{1}{2}(1 \pm \gamma_5)$ and e_q is the quark charge in units of $e > 0$. (For gluino interactions, replace $e_q e$ with $g_s T^a$.) The

* If Goldstinos exist, then we should also consider the decay $\tilde{g} \to \tilde{G} + q\bar{q}$. This decay mode has a lifetime [5] $\tau = 192\pi^2 M_s^4/\alpha_s \tilde{M}_g^5$, and a signature identical to that of $\tilde{g} \to \tilde{\gamma} + q\bar{q}$. As discussed in the text, if M_s is large then this mode will also be negligible.

computation of the decay rate of $\tilde{g}$ is analogous to that of μ-decay; we give the result only for massless quarks and $\tilde{\gamma}$ (The correction for massive $\tilde{\gamma}$ can be obtained by integrating eq. (4) over x.):

$$\Gamma(\tilde{g} \to \tilde{\gamma} q \bar{q}) = \frac{\alpha_s \alpha e_q^2 \tilde{M}_g^5}{96\pi} \left(\frac{1}{\tilde{m}_L^4} + \frac{1}{\tilde{m}_R^4} \right). \tag{3}$$

The photino energy spectrum is also a useful quantity as it describes the distribution of missing energy in the gluino decay. Writing $x = E_{\tilde{\gamma}}/\tilde{M}_g$, and $\tilde{x} = m_{\tilde{\gamma}}/\tilde{M}_g$, the photino energy distribution including the possibility of a non-zero photino mass, is

$$\mathrm{d}\Gamma/\mathrm{d}x \sim (x^2 - \tilde{x}^2)^{1/2} \left[x(1 - \tfrac{4}{3}x) + \tilde{x}(1 - 2x) - (\tfrac{2}{3} - x)\tilde{x}^2 + \tilde{x}^3 \right]. \tag{4}$$

The range for x is $\tilde{x} \leqslant x \leqslant \frac{1}{2}(1 + \tilde{x}^2)$. Note this includes constructive interference effects between figs. 1a and 1b.

There are three basic ways to detect the gluino. They are useful because the gluino production cross section is large, since gluinos are color octet states; the cross section is about an order of magnitude larger than for a heavy quark of the same mass [3]. First, if the gluino mass is light enough, then it is possible [2–5] that there exist long-lived charged or neutral supersymmetric hadrons which leave visible tracks or gaps. Second, the gluino might be discovered by seeing the emitted photino interact [2, 3, 16]. This is the idea behind beam dump searches [17] where the photino events would be interpreted as excess neutrino induced neutral current events. The non-observation of the gluino due to the above two methods can at present set limits for the gluino mass of roughly $\tilde{M}_g \gtrsim 2$–5 GeV, depending on the scalar quark masses since these affect the photino interaction rate. Masses measured or limited this way include constituent mass, presumably of order one GeV. The third technique for discovering the gluino uses hadron colliders where gluinos are pair produced or singly produced with either a photino or a scalar quark. When the gluinos decay the final state photinos, which interact rather weakly, escape the collider detector, and provide events with missing energy and momentum but with no hard charged lepton.

In the gluino pair events just described, the topology of the event depends crucially on whether the gluino decays into two or three bodies. Assuming the goldstino modes are irrelevant, the dominant decay of the gluino had been considered to be into three bodies, $q\bar{q}\tilde{\gamma}$. Our main point in this paper is that there is a two-body decay mode which has not previously been discussed in the literature, $\tilde{g} \to g\tilde{\gamma}$. This decay violates C-invariance, so naively it appears that it should be negligible. However, in supersymmetric theories, it is natural to have different masses for the $\tilde{q}_L$ and $\tilde{q}_R$ (the mass splitting is due to weak interaction effects) which we denote by $\tilde{m}_L$ and $\tilde{m}_R$. This will lead to parity (P) and (assuming CP) C-non-invariance in the strong interactions. The non-observation of such effects in

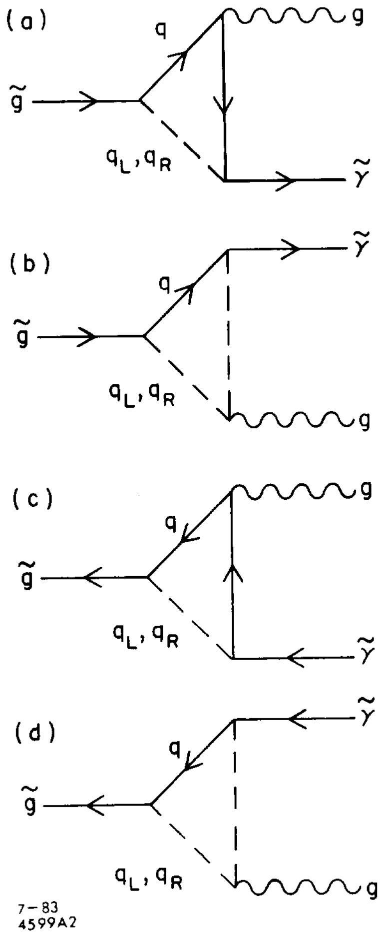

Fig. 2. The $\tilde{g} \to g\tilde{\gamma}$ decay. Note that diagrams (a) and (c) are identical except for the direction of the arrow inside the loop (similarly for (b) and (d)). The direction of the arrows on the external Majorana fermion lines ($\tilde{\gamma}$ and $\tilde{g}$) is arbitrary and chosen as shown for convenience. One must sum over all possible flavors in the loop.

ordinary strong interactions restricts $\tilde{m}_L - \tilde{m}_R$ for scalars partners of the u and d quarks as shown in refs. [18, 19], but there are no restrictions for second or third generation scalar quarks. It is quite plausible that the decay rate for $\tilde{g} \to g\tilde{\gamma}$ will be substantial.

The result for $\Gamma(\tilde{g} \to g\tilde{\gamma})$ is calculated in sect. 3, but we give it here for continuity. The Feynman graphs for the decay are shown in fig. 2. If we assume that the gluino and quark masses are small compared to $\tilde{m}_L$ and $\tilde{m}_R$,

$$\Gamma(\tilde{g} \to g\tilde{\gamma}) = \frac{\alpha_s^2 \alpha e_q^2 \tilde{M}_g^5}{128\pi^2} \sum_{\text{flavors}} \left(\frac{1}{\tilde{m}_L^2} - \frac{1}{\tilde{m}_R^2} \right)^2 . \tag{5}$$

The sum in eq. (5) is over all flavors of scalar quarks (which appear in the loops of fig. 2) and the $\tilde{\gamma}$ mass has been set to zero. Note that an advantage which could favor $\tilde{g} \to g\tilde{\gamma}$ over $\tilde{g} \to \tilde{\gamma}q\bar{q}$ is that in the former case all quark flavors contribute (although if $m_q \neq 0$ the formula is modified such that the contribution from heavy quarks is suppressed; see eq. (14)), whereas in the latter case only quarks permitted by kinematics are allowed (with possible phase space suppression). In order to compare eqs. (3) and (5), let us count only one flavor of quarks. Then,

$$\frac{\Gamma(\tilde{g} \to g\tilde{\gamma})}{\Gamma(\tilde{g} \to q\bar{q}\tilde{\gamma})} \approx \frac{3\alpha_s\left(\tilde{m}_R^2 - \tilde{m}_L^2\right)^2}{4\pi\left(\tilde{m}_L^4 + \tilde{m}_R^4\right)}. \tag{6}$$

The strong running coupling constant is evaluated at $Q^2 = \tilde{M}_g^2$. Thus, taking $\alpha_s \approx \frac{1}{5}$ and accounting for the other quark flavors we find that the suppression of the $g\tilde{\gamma}$ decay mode relative to the $\bar{q}q\gamma$ mode depends on the size of the factor $(\tilde{m}_R^2 - \tilde{m}_L^2)^2/(\tilde{m}_L^4 + \tilde{m}_R^4)$. How large is this factor? Large splittings between the masses of the scalar partners of the u and d quarks could lead to the appearance of P-violation in the strong interactions. Limits on this mass splitting for up and down scalar quarks have been given in refs. [18,19]; they depend on the masses of the gluinos and scalar quarks. The larger these masses are, the less restrictive the bounds on the mass splittings. For scalar $\tilde{u}$ and $\tilde{d}$ quark masses above 100 GeV, this ratio can be as large as 1. There are no restrictions on this ratio for s, c, b, t quark partners.

It is interesting that one is seeing here a parity violation in the strong interactions. Of course, the effect is generated by having $\tilde{m}_{q_L} \neq \tilde{m}_{q_R}$, which in turn is induced by electroweak effects, so the fundamental violation is still rooted in the expected place. Presumably there are a number of other places to look for parity and charge-conjugation violation in high-energy interactions and heavier flavor production.

Various models of low-energy broken supergravity can give us some idea what values one might expect for the mass splitting. For example, ref. [14] predicts:

$$\tilde{m}_{u_L}^2 - \tilde{m}_{u_R}^2 = \left(\tfrac{4}{3}\sin^2\theta_W - \tfrac{1}{2}\right)m_Z^2, \tag{7a}$$

$$\tilde{m}_{d_L}^2 - \tilde{m}_{d_R}^2 = \left(-\tfrac{2}{3}\sin^2\theta_W + \tfrac{1}{2}\right)m_Z^2, \tag{7b}$$

$$\tilde{m}_{q_L}^2 + \tilde{m}_{q_R}^2 = 2m_{g_{3/2}}^2 \mp \tfrac{1}{2}m_Z^2, \qquad \text{for } q = \begin{pmatrix} u \\ d \end{pmatrix}, \tag{7c}$$

($m_Z \approx 95$ GeV), for the light up (u) and down (d) type quarks. Thus, depending on the exact value of $m_{g_{3/2}}$, it seems possible that the rate for $\tilde{g} \to g\tilde{\gamma}$ is non-negligible. Other models (e.g. Ellis, et al., ref. [12]) have significant $\tilde{q}_L - \tilde{q}_R$ mass splitting only in the heavy (top) quark sector. In this case, too, we find a two-body decay rate which could be competitive with the three-body modes. Implications for phenomenology will be discussed in sect. 4.

So far, we have assumed that the $\tilde{\gamma}$ is the lightest supersymmetric particle. If there are other supersymmetric particles lighter than the $\tilde{g}$, then new decay channels

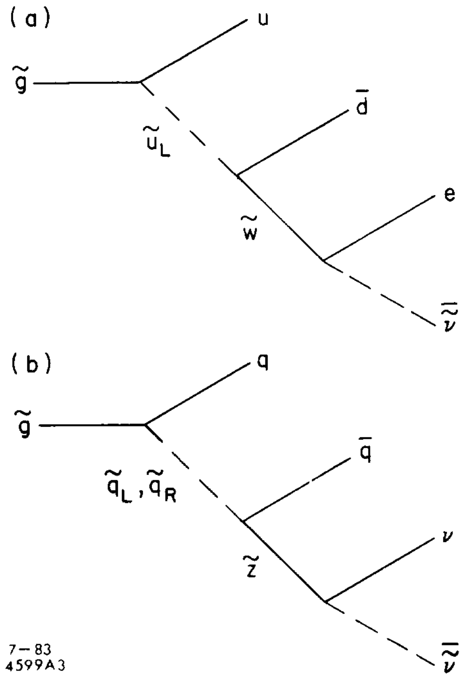

Fig. 3. Four-body decays of the gluino. If the scalar neutrino ($\tilde{\nu}$) is lighter than the gluino, then these diagrams are allowed. In addition to the diagrams shown, we must also include diagrams whose final states are the charged conjugates to the ones shown and diagrams with other flavors of quarks, leptons, and scalar neutrinos.

become available. One could imagine situations [15] where the $\tilde{\gamma}$ and $\tilde{g}$ were nearly degenerate so that the usual $\tilde{g}$ decay modes would be phase-space suppressed. Alternatively, one might use cosmological arguments to rule out a light stable photino [20]. In these cases, the existence of alternative decay modes could be important. The main candidate for another light supersymmetric particle is the scalar neutrino ($\tilde{\nu}$) [7]*. Although various models previously referred to prefer a heavy scalar neutrino, there is at present no experimental evidence against a light scalar neutrino [8, 21, 22].

A light $\tilde{\nu}$ would allow for four-body decay modes of the gluino (see fig. 3). Their rates are calculated in sect. 5; the result is given here. Taking all final state particles massless and all internal lines superheavy, we find (for $\tilde{M}_g \ll m_W, \tilde{m}_L$):

$$\Gamma(\tilde{g} \to u\bar{d}e^-\bar{\tilde{\nu}} + \text{c.c.}) = \frac{\alpha^2 \alpha_s \tilde{M}_g^9}{32 \times 45 \times 64\pi^2 \sin^4\theta_w \tilde{m}_L^4 m_W^4}, \tag{8}$$

* This possibility was independently noticed in ref. [8].

where c.c. stands for charge conjugate states. Adding to this the contribution of fig. 3b and allowing for additional flavors of quarks and leptons in the final state would multiply the result of eq. (8) by roughly a factor of 10. The results here are analogous to those of ref. [22]. If the photino is light so that eqs. (3) and (5) are relevant, then the four-body decays are negligible for $\tilde{M}_g \ll \tilde{m}_L, m_W$. On the other hand, if the two- and three-body decays happen to be kinematically suppressed, or for heavy gluino masses (where the approximation which led to eq. (8) is no longer valid), the four-body decays could be appreciable. These would lead to interesting signatures as discussed in sect. 5.

3. Calculation of $\tilde{g} \to g\tilde{\gamma}$

We now turn to details of the calculation of $\Gamma(\tilde{g} \to g\tilde{\gamma})$. The first observation is that $\tilde{g}$ and $\tilde{\gamma}$ are Majorana fermionic partners of the gluon and photon respectively. Hence both the $\tilde{g}$ and $\tilde{\gamma}$ have charge conjugation number -1 (in the absence of C-violating interactions) [23]. It follows that $\tilde{g} \to g\tilde{\gamma}$ is forbidden unless C-invariance is violated (although CP-invariance is still maintained). The matrix element must then have the following form: [24]

$$M = \frac{iA}{\tilde{M}_g} \bar{u}(k_1)\gamma_5\sigma_{\mu\nu}u(p)\varepsilon^*_\mu(k_2)k_{2\nu}, \tag{9}$$

where the momenta are as follows: $\tilde{g}(p) \to \tilde{\gamma}(k_1) + g(k_2)$, $\tilde{M}_g$ is the gluino mass and ε_μ is the gluon polarization vector. Eq. (9) follows from the requirements of CP invariance and gauge invariance*. We may easily compute the rate by averaging over initial colors and spin and including the color factor: $\mathrm{Tr}(T^aT^b)\mathrm{Tr}(T^aT^b) = \frac{1}{4}\delta^{ab}\delta^{ab} = 2$. The result is:

$$\Gamma(\tilde{g} \to g\tilde{\gamma}) = \frac{|A|^2}{32\pi} \frac{\left(\tilde{M}_g^2 - \tilde{m}_\gamma^2\right)^3}{\tilde{M}_g^5}. \tag{10}$$

The $\tilde{g} \to g\tilde{\gamma}$ decay occurs via the one-loop diagrams depicted in fig. 2.

For simplicity, we take $\tilde{q}_L$ and $\tilde{q}_R$ to be the appropriate scalar quark eigenstates with masses $\tilde{m}_L$ and $\tilde{m}_R$ respectively. The relevant supersymmetric interaction terms are given by eq. (2) and below [25]:

$$ig_s T^a\left[\tilde{q}_L^* \overleftrightarrow{\partial}_\mu \tilde{q}_L + \tilde{q}_R^* \overleftrightarrow{\partial}_\mu \tilde{q}_R\right] A_a^\mu. \tag{11}$$

Note that in fig. 2 the graphs (a) and (c) are identical except for the direction of the

* This form was found by Pal and Wolfenstein (ref. [24]) when they computed radiative decays of Majorana neutrinos, $\nu_2 \to \nu_1 + \gamma$, where ν_1 and ν_2 had the same CP properties.

arrow inside the loop (similarly for (b) and (d)*). This is analogous to the situation in QED when one proves Furry's Theorem [27] perturbatively. There the two graphs which are identical, except for the direction of the arrows inside the loop, exactly cancel thus causing the three-photon Green function to vanish. This is simply a consequence of C-invariance in QED. In our case, if $\tilde{m}_{\mathrm{L}} = \tilde{m}_{\mathrm{R}}$, a similar phenomenon occurs – the sum of all the graphs of fig. 2 vanishes. If $\tilde{m}_{\mathrm{L}} \neq \tilde{m}_{\mathrm{R}}$, the sum leads to a non-vanishing result.

The calculation is straightforward and we only quote the final results. Note that individual graphs are divergent but the infinities cancel when all graphs are summed. We indeed find that the final amplitude is of the form given by eq. (9). The result for A is**

$$A = \frac{g_s^2 e e_{\mathrm{q}}}{4\pi^2}\left(I_{\mathrm{L}} - I_{\mathrm{R}}\right), \tag{12}$$

where

$$I_{\mathrm{i}} = 1 + \frac{1}{\tilde{M}_{\mathrm{g}}^2}\int_0^1 \frac{\mathrm{d}x}{x(1-x)}\left[M_{\mathrm{q}}^2 x + \tilde{m}_{\mathrm{i}}^2(1-x)\right] \times \log\left(1 - \frac{\tilde{M}_{\mathrm{g}}^2 x(1-x)}{m_{\mathrm{q}}^2 x + \tilde{m}_{\mathrm{i}}^2(1-x)}\right). \tag{13}$$

I_{i} may be expressed in terms of dilogarithms; the exact expressions are given in the appendix. Here we note some useful limits: let $R \equiv \tilde{M}_{\mathrm{g}}^2/\tilde{m}_{\mathrm{i}}^2$ and $r \equiv m_{\mathrm{q}}^2/\tilde{m}_{\mathrm{i}}^2$. For small gluino mass, ($\tilde{M}_{\mathrm{g}} \ll \tilde{m}_{\mathrm{i}}$),

$$I_{\mathrm{i}} = \frac{-R}{4(1-r)^2}\left[1 + r + \frac{2r\log r}{1-r}\right] + \mathrm{O}(R^2). \tag{14}$$

Note that I_{i} vanishes in the limit of zero gluino mass. Second, for light quarks, ($r = 0$) we obtain the result:

$$I_{\mathrm{i}} = 1 - \frac{\mathrm{Li}_2(R)}{R} + \mathrm{O}(r). \tag{15}$$

Finally, if $m_{\mathrm{q}} = \tilde{m}_{\mathrm{i}}$, then we find

$$I_{\mathrm{i}} = 1 - \frac{4}{R}\left[\sin^{-1}\left(\tfrac{1}{2}\sqrt{R}\right)\right]^2. \tag{16}$$

* The direction of the arrows on the external Majorana fermion lines is arbitrary. The choice made in fig. 2 allows us to use simpler Feynman rules. See the appendix of ref. [26].

** This occurs as follows. The sum of diagrams a and b of fig. 2 gives a matrix element proportional to $\bar{u}(P_{\mathrm{L}} I_{\mathrm{L}} + P_{\mathrm{R}} I_{\mathrm{R}})\sigma_{\mu\nu} u$. From this, the sum of diagrams c and d can be obtained by interchanging P_{L} and P_{R} and multiplying by -1.

The total rate is obtained by summing over all flavors j,

$$\Gamma(\tilde{g} \to g\tilde{\gamma}) = \frac{\alpha_s^2 \alpha e_Q^2}{8\pi^2} \frac{\left(\tilde{M}_g^2 - \tilde{m}_\gamma^2\right)^3}{\tilde{M}_g^5} \sum_j \left(I_{L_j} - I_{R_j}\right)^2 . \tag{17}$$

In the approximation that both the gluino and internal quarks are light, either eq. (14) or eq. (15) yield $I_i \approx -\frac{1}{4}R$ which immediately leads to eq. (5) (when $\tilde{m}_\gamma = 0$).

4. Signatures of gluinos

We have argued in sect. 2 that the $g\tilde{\gamma}$ mode may be competitive with the $q\bar{q}\tilde{\gamma}$ mode. This observation has important consequences in searching for gluinos produced in hadronic colliders or in beam dump experiments. First, consider the missing energy carried off by $\tilde{\gamma}$ from a gluino decay at rest. The $\tilde{\gamma}$ energy from the two-body decay is simply $E_{\tilde{\gamma}} = (\tilde{M}_g^2 + \tilde{m}_\gamma^2)/2\tilde{M}_g \geqslant 0.5\ \tilde{M}_g$ in the rest frame of the decaying gluino (for zero mass gluons), whereas this energy is the endpoint of a distribution in the case of a three-body decay. The average $\tilde{\gamma}$ energy from the three-body decays is $E = 0.35\ \tilde{M}_g$ for massless photinos; it stays at that value for $\tilde{x} = \tilde{m}_\gamma/\tilde{M}_g \lesssim 0.15$, and then increases almost linearly to unity when $\tilde{x} = 1$. This average energy is always less than $E_{\tilde{\gamma}}$ for the two-body mode. Furthermore, an important consideration in searching for gluinos in beam dump experiments, where the final photino is detected, is how much p_T is induced by the decay; the greater the p_T, the fewer final photinos will end up in the detector downstream. As is well known, three-body decays of heavy objects typically induce a significant p_T, while for two-body decays of a moving particle almost all decay products emerge near the minimum opening angle given by $\cos\frac{1}{2}\theta = v$, where v is the velocity of the decaying particle. Thus the two-body decay will improve the sensitivity of beam dump searches for gluinos; the precise effect has to be estimated for each experimental situation with the geometry and detection efficiencies taken into account.

For hadronic colliders the $g\tilde{\gamma}$ decay mode will always produce a considerably cleaner signal. Suppose a pair of $\tilde{g}$ is produced, and both gluinos decay $\tilde{g} \to g\tilde{\gamma}$. The event will look like an acoplanar two-jet event with substantial missing energy, and it may be possible to detect a signal on an event-by-event basis, rather than by applying appropriate cuts to the data on a statistical basis (see e.g. ref. [28]). There is very little standard model background for this signature if hard charged leptons can be vetoed with confidence. Since in practice both $g\tilde{\gamma}$ and $q\bar{q}\tilde{\gamma}$ final states may occur, one can expect combination events with different decays from different $\tilde{g}$.

A rarer process which might lead to even more spectacular signals and hence be easily observable* is $q\bar{q} \to \tilde{g}\tilde{\gamma}$ which has the potential for producing "one-jet" events

* This has been advocated by A. Savoy-Navarro as a very good way to search for gluino production.

with no accompanying leptons. Again, the decay $\tilde{g} \to g\tilde{\gamma}$ would lead to a final state (at large p_T) of $g\tilde{\gamma}\tilde{\gamma}$ which seems like quite an attractive signature for supersymmetry. Production of a Z^0 which decays into a neutrino pair and is accompanied by a gluon jet will be a background for these events. Detailed calculations using a Monte Carlo simulation of hadronic collisions (e.g. ISAJET [29]) will be necessary to determine more precisely how well one can observe gluinos by these methods.

To compare the three- and two-body decays, one further point is worth noting. Consider the case where the $\tilde{\gamma}$ mass is significant (i.e. $\tilde{x}$ near 1). From eqs. (4) and (10), it follows that the three-body rate vanishes as $(1-\tilde{x})^2$ relative to the two-body rate (this is simply a consequence of the well-known $(\Delta m)^5$ factor which occurs in β-decay). Hence, if $\tilde{x} \gtrsim \frac{1}{2}$, the two-body mode is significantly enhanced relative to the three-body decay. Exact formulas for the three-body decay rate for arbitrary masses ($\tilde{x}$) will be given by us in ref. [1].

5. Four-body decay

If there exist other supersymmetric particles lighter than the gluino, then new $\tilde{g}$ decay modes will be possible. In particular, the scalar neutrino ($\tilde{\nu}$) could be the lightest supersymmetric particle [7]. At present, there are no strong experimental limitations on the mass of the $\tilde{\nu}$ [8, 22]; the only constraint [21] is that $\tilde{m}_\nu$ for $\tilde{\nu}_\tau$, $\tilde{\nu}_\mu$, $\tilde{\nu}_e$ must be large enough so the decay $\tau \to \tilde{\nu}_\tau \bar{\tilde{\nu}}_\ell \ell$ is not observable. Especially if photinos are heavy, it may be kinematically favored for the gluino to decay into a $\tilde{\nu}$ and light quarks and leptons. Possible decay modes are depicted in fig. 3. If the emitted $\tilde{\nu}$ is stable (or if the dominant $\tilde{\nu}$ decay is into $\nu\tilde{\gamma}$) then the $\tilde{\nu}$ four-momentum will be experimentally unobservable [22]. The computation of the gluino four-body decay rates is analogous to the four-body decay rates of the $\tilde{\nu}$ of ref. [22], so we will simply outline the procedure here.

For illustration purposes, we compute the diagram in fig. 3a: $\tilde{g}(p) \to u(k_1) + \bar{d}(k_2) + e^-(k_3) + \bar{\tilde{\nu}}(k_4)$ (the four-momenta are given in parentheses). The relevant Feynman rules are given in ref. [30]. The matrix element squared, summed and averaged over colors and spins is:

$$|M|^2 = \frac{256\pi^3\alpha^2\alpha_s}{m_W^4 \tilde{m}_L^4 \sin^4\theta_w} p\cdot k_1 k_2\cdot k_4 k_3\cdot k_4, \tag{18}$$

where we have assumed that $\tilde{M}_g \ll m_W, \tilde{m}_L$ and all final state particles are taken to be massless. Now integrate over massless four-body phase space. In the rest frame of the gluino,

$$\Gamma = \frac{\alpha^2\alpha_s}{24\pi^2 m_W^4 \tilde{m}_L^4 \sin^4\theta_w} \int E_3\, dE_3\, E_4\, dE_4\, dx\, G(E_3, E_4, x), \tag{19}$$

where $E_j \equiv |\boldsymbol{k}_j|$ $(j=3,4)$, $x \equiv \hat{\boldsymbol{k}}_3 \cdot \hat{\boldsymbol{k}}_4$ and

$$G(E_3, E_4, x) = E_3 E_4^2 \Big[(1-x) \tilde{M}_g \big[4(E_3+E_4) - 3\tilde{M}_g \big] + 2E_3 \big(\tilde{M}_g - E_3 - 2E_4 \big)(1-x)^2 \Big] . \tag{20}$$

The limits of integration can be obtained from ref. [31]:

(i) $0 \leqslant E_4 \leqslant \frac{1}{2}\tilde{M}_g$,

(ii) $0 \leqslant E_3 \leqslant \frac{1}{2}\tilde{M}_g$,

(iii) If $0 \leqslant E_3 \leqslant \frac{1}{2}\tilde{M}_g - E_4$, then $-1 \leqslant x \leqslant 1$,

(iv) If $\frac{1}{2}\tilde{M}_g - E_4 \leqslant E_3 \leqslant \frac{1}{2}\tilde{M}_g$, then

$$-1 \leqslant x \leqslant 1 - \left(\frac{2\tilde{M}_g(E_3+E_4) - \tilde{M}_g^2}{2E_3E_4} \right).$$

The integration is straightforward and tedious; the final result is given in eq. (8) after adding the charge conjugated final state. We conclude that if the $\tilde{\gamma}$ and $\tilde{\nu}$ are light, then the four-body decay branching ratios are completely negligible unless $\tilde{M}_g$ and $\tilde{m}_L$ are nearly equal to m_W. In such a case, the total four-body decays (summing all diagrams in fig. 3) although somewhat rare will be non-negligible if gluinos are copiously produced. The $\tilde{g} \to q\bar{q}\nu\bar{\tilde{\nu}}$ mode will look similar to $\tilde{g} \to q\bar{q}\tilde{\gamma}$. However, $\tilde{g} \to u\bar{d}e^-\bar{\tilde{\nu}}$ is unique in that it is the only gluino decay mode which contains a primary electron. (This channel could be confused with $\tilde{g} \to c\bar{s}\tilde{\gamma}$ where the c-quark decays semileptonically.)

The four-body decay modes are complicated in structure and are likely to confuse the issue rather than elucidate the possible existence of gluinos. However, one should be aware that in unusual circumstances where the $\tilde{\gamma}$ is heavy but the $\tilde{\nu}$ is light, the four-body decays could be the dominant ones. If this should happen, the lifetime associated with eq. (8) is quite long ($\tau \sim 10/\tilde{M}_g^9$ sec. with $\tilde{M}_g$ in GeV) in which case gluinos would show up in long-lived heavy particles.

6. Summary

In this paper we have studied various possible decay modes of the gluino. This is of particular interest to experimentalists who need to know the signatures of gluinos produced in their experiments. Gluinos, if they exist, are expected to be copiously produced in hadronic colliders as long as the c.m. energy is sufficiently greater than the gluino mass. In most supersymmetric models, it had been expected that $\tilde{g} \to q\bar{q}\tilde{\gamma}$ was the dominant gluino decay mode. We have shown in this paper that the C-violating two-body decay mode $\tilde{g} \to g\tilde{\gamma}$ could be competitive with the three-body decay modes. If $\tilde{g} \to g\tilde{\gamma}$, the gluino signature at hadronic colliders and in beam dump experiments is likely to be considerably cleaner. In such a case, if gluinos are produced via $gg \to \tilde{g}\tilde{g}$, the event will appear as two acoplanar jets. Also dramatic will be $q\bar{q} \to \tilde{g}\tilde{\gamma}$ which will appear as one gluon jet and a lot of missing momentum and energy (with no accompanying hard charged leptons). With one or both of these production mechanisms and the $g\tilde{\gamma}$ decay it may be possible to see $\tilde{g}$ production on an individual event basis rather than statistically. We have also studied the possibility of four-body decays which can occur if the scalar neutrino is lighter than the gluino, and could dominate if $\tilde{m}_\gamma \simeq \tilde{M}_g$. These four-body branching ratios are suppressed except in such exceptional conditions.

Unless they are extremely heavy, gluinos are probably the easiest supersymmetric particle to produce. Their detection will depend on the decay rates discussed in this paper. In the next few years experiments at the ISR, SPS, the fixed target Tevatron, and the Tevatron Collider all will be sensitive to finding gluinos. If nature were supersymmetric, the gluino could be the first window to physics beyond the standard model.

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Appendix

Here, we evaluate the integral I_i given in eq. (13). Using the notation $R \equiv \tilde{M}_g^2/\tilde{m}_i^2$ and $r \equiv m_q^2/\tilde{m}_i^2$, we must evaluate

$$I = f(R, r) + f(R/r, 1/r) + 1, \tag{A.1}$$

where

$$f(R, r) = \frac{1}{R}\int_0^1 \frac{dx}{x} \log\left(1 - \frac{Rx(1-x)}{1-x+rx}\right). \tag{A.2}$$

Note that if $\tilde{M}_g < m_q + \tilde{m}_i$ (i.e. $0 \leqslant R \leqslant 1 + r + 2\sqrt{r}$) then $f(R, r)$ is a real valued function. Using the methods of Lewin [32],

$$f(R, r) = \frac{1}{R}\left\{\mathrm{Li}_2(1-r) - \mathrm{Li}_2\left[\tfrac{1}{2}(1+R-r+\lambda)\right] - \mathrm{Li}_2\left[\tfrac{1}{2}(1+R-r-\lambda)\right]\right\}, \tag{A.3}$$

where $\lambda \equiv [(1-R-r)^2 - 4Rr]^{1/2}$ and the dilogarithm is defined as

$$\mathrm{Li}_2(x) = -\int_0^x \frac{\log(1-t)\,\mathrm{d}t}{t}. \tag{A.4}$$

For $1 + r - 2\sqrt{r} \leqslant R \leqslant 1 + r + 2\sqrt{r}$, the following expression in place of (A.3) is more useful:

$$f(R,r) = \frac{1}{R}\left[\mathrm{Li}_2(1-r) - 2\,\mathrm{Li}_2(\sqrt{R}, \theta)\right], \tag{A.5}$$

where $\cos\theta \equiv (1 + R - r)/(2\sqrt{R})$, $0 \leqslant \theta \leqslant \pi$, and

$$\mathrm{Li}_2(x, \theta) = -\tfrac{1}{2}\int_0^x \frac{\log(t^2 - 2t\cos\theta + 1)\,\mathrm{d}t}{t} \tag{A.6}$$

or equivalently, $\mathrm{Li}_2(x,\theta) = \mathrm{Re}\,\mathrm{Li}_2(x\,\mathrm{e}^{i\theta})$ for x real.

The properties [32] of the functions $\mathrm{Li}_2(x)$ and $\mathrm{Li}_2(x,\theta)$ allow us to deduce the following interesting special cases.

(i) $R \to 0$ limit.

$$f(R,r) = \frac{-1}{1-r}\left[1 + \frac{r\log r}{1-r}\right] - \frac{R}{2(1-r)^2}\left[\tfrac{1}{2}(1+5r) + \frac{r(2+r)\log r}{1-r}\right] + \mathrm{O}(R^2). \tag{A.7}$$

(ii) $r \to 0$ limit.

$$f(R,r) = \frac{-\mathrm{Li}_2(R)}{R} + \mathrm{O}(r). \tag{A.8}$$

(iii) $r = 1$.

$$f(R,1) = -\frac{2}{R}\left[\sin^{-1}\left(\tfrac{1}{2}\sqrt{R}\right)\right]^2, \tag{A.9}$$

where

$$-\tfrac{1}{2}\pi \leqslant \sin^{-1}\left(\tfrac{1}{2}\sqrt{R}\right) \leqslant \tfrac{1}{2}\pi.$$

Note that the above formulas were derived under the assumption that $R \leqslant 1 + r + 2\sqrt{r}$. (For R above this value, $f(R,r)$ is complex. Its analytic continuation can be performed using eq. (A.3). However, extreme care must be taken with the cut structure. Note that $\mathrm{Li}_2(x)$ is analytic in the complex plane cut from $x = 1$ to ∞.)

Using eqs. (A.1), (A.7)–(A.9), the eqs. (14)–(16) in the text immediately follow.

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